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Neutrino texture saturating the CP asymmetry

Takayasu Baba and Daijiro Suematsu [†]

Institute for Theoretical Physics, Kanazawa University,

Kanazawa 920-1192, Japan

Abstract

We study a neutrino mass texture which can explain the neutrino oscillation data and also saturate the upper bound of the CP asymmetry ε_1 in the leptogenesis. We consider the thermal and non-thermal leptogenesis based on the right-handed neutrino decay in this model. A lower bound of the reheating temperature required for the explanation of the baryon number asymmetry is estimated as $O(10^8)$GeV for the thermal leptogenesis and $O(10^6)$GeV for the non-thermal one. It can be lower than the upper bound of the reheating temperature imposed by the cosmological gravitino problem. An example of the construction of the discussed texture is also presented.

[†]e-mail: suematsu@hep.s.kanazawa-u.ac.jp

1 Introduction

The discovery of the neutrino masses [1] gives large impact to the study of particle physics and astroparticle physics. In particular, it presents an interesting approach to the study of the origin of the baryon number (B) asymmetry in the universe, which is one of the important questions in these fields. The leptogenesis [2] based on the $B-L$ violation due to the neutrino masses is considered to be the most promising scenario for the generation of the B asymmetry. During the recent few years, the leptogenesis based on the CP asymmetric decay of the heavy right-handed neutrinos [3] whose existence is required by the seesaw mechanism [4] has been extensively studied [5, 6].

Since the intermediate scale is generally necessary for the seesaw mechanism, it seems to be natural to consider the leptogenesis in the supersymmetric framework to guarantee the stability of that scale against the radiative correction. However, if we consider it in such a framework, a crucial problem called the cosmological gravitino problem is caused in relation to the generation of the right-handed neutrinos. If the reheating temperature T_R required to produce a sufficient amount of the right-handed neutrinos is a high value such as $T_R \gtrsim 10^8$GeV, the gravitino can be produced too much and its late time decay may disturb the nucleosynthesis [7].[1] Thus, the production mechanism of the heavy right-handed neutrinos is the important ingredient for this problem. Several solutions for this difficulty have been proposed by now [9, 10].

On the other hand, the CP asymmetry [3] in the decay of the right-handed neutrinos is another crucial factor which plays the essential role to determine the generated lepton number (L) asymmetry. Its magnitude depends on the structure of the neutrino mass matrix, which is severely constrained to explain the neutrino oscillation data [1]. From a viewpoint of the leptogenesis, it is favorable that the neutrino mass matrix can realize the maximum value of the CP asymmetry [3, 11] or enhance its value [12, 13]. Thus it is important for the quantitative study of the leptogenesis to construct such a concrete model for the neutrino mass matrix as done in [5, 14, 15] and to proceed the investigation based on it. In this paper we present an example of the neutrino mass matrix and estimate the reheating temperature required to produce the sufficient B asymmetry.

The paper is organized as follows. In section 2 we present a neutrino mass texture

[1] If the gravitino is the lightest superparticle as in the gauge mediated supersymmetry breaking, there is no gravitino problem even in the case of $T_R = O(10^{10})$GeV [8].

and discuss its phenomenological features. An example for its construction is discussed in appendix A. In section 3 we apply this model to both the thermal and non-thermal leptogenesis. We discuss the lower bound of the reheating temperature required for the production of the sufficient B asymmetry. Section 4 is devoted to the summary.

2 Neutrino mass texture

We consider the minimal supersymmetric standard model (MSSM) extended with gauge singlet chiral superfields N_i which correspond to three generation right-handed neutrinos. An effective superpotential for the neutrino sector is assumed as follows:

$$W = \sum_{i,j=1}^{3} \left(h^{\nu}_{ij} N_i H_2 L_j + \frac{1}{2} \mathcal{M}_{ij} N_i N_j \right), \tag{1}$$

where L_i and H_2 are the lepton doublet and Higgs doublet chiral superfields, respectively. In this paper we use the same notation for both a superfield and its component fields. The right-handed neutrino mass matrix $\mathcal{M}$ and the Dirac mass matrix m_D induced from the first term in this superpotential are assumed to take the form as[2]

$$\mathcal{M} \equiv \begin{pmatrix} M_1 & m & 0 \\ m & M_2 & 0 \\ 0 & 0 & M_3 \end{pmatrix}, \qquad m_D \equiv h^{\nu} \langle H_2 \rangle = \begin{pmatrix} 0 & 0 & 0 \\ a & a' & 0 \\ 0 & b & b' \end{pmatrix}, \tag{2}$$

where the charged lepton mass matrix is considered to be diagonal. Although each element of $\mathcal{M}$ is supposed to be real, m_D is assumed to be a complex matrix.

If the hierarchical structure

$$m, \ M_1 \ll M_2 \ll M_3 \tag{3}$$

is assumed in the right-handed neutrino sector, the eigenvalues $\tilde{M}_i$ of the mass matrix $\mathcal{M}$ can be approximated to

$$\begin{array}{llll} \text{(a)} & \tilde{M}_1 (\simeq M_1), & M_2, \quad M_3, & (\text{for } m^2 < M_1 M_2), \\ \text{(b)} & \tilde{M}_1 (\simeq M_2 \sin^2 \xi), & M_2, \quad M_3, & (\text{for } m^2 > M_1 M_2), \end{array} \tag{4}$$

[2] This model can be considered as a simple extension of the model in [16]. It is realized by adding a right-handed neutrino N_1 to the original one in such a way that N_1 weakly couples to N_2 alone.

where we use $\tilde{M}_i \simeq M_i$ $(i = 2, 3)$ and $\sin\xi \simeq m/M_2$. These two cases are studied in the following part. The structure of $\mathcal{M}$ and m_D can be effectively realized by imposing a suitable symmetry on the superpotential at the high energy scales. We give such an example for the construction of $\mathcal{M}$ and m_D in appendix A.

If we change N_i into the $\mathcal{M}$ diagonal basis $\tilde{N}_i$, the Dirac neutrino mass matrix is transformed into

$$\tilde{m}_D \equiv \tilde{h}^\nu \langle H_2 \rangle = \begin{pmatrix} -a\sin\xi & -a'\sin\xi & 0 \\ a\cos\xi & a'\cos\xi & 0 \\ 0 & b & b' \end{pmatrix}. \tag{5}$$

Applying the seesaw mechanism to these matrices, we can obtain the light neutrino mass eigenvalues and the MNS matrix. Here, for the simplicity, we put $a' = \sqrt{2}a$ and $b' = b$,[3] and then the light neutrino mass eigenvalues are found to be

$$m_1 = 0, \qquad m_2 \simeq \frac{2|a|^2}{M_2}, \qquad m_3 \simeq \frac{2|b|^2}{M_3}. \tag{6}$$

The MNS matrix has the bi-large mixing form such as

$$U \simeq \begin{pmatrix} \frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}}\sin\theta \\ -\frac{1}{2} & \frac{1}{2}\cos\theta & \frac{1}{\sqrt{2}}\cos\theta \\ \frac{1}{2} & -\frac{1}{2}\cos\theta & \frac{1}{\sqrt{2}}\cos\theta \end{pmatrix}, \tag{7}$$

where we neglect the CP phase in this expression.

Now we can compare these results with the present experimental data. Since the neutrino mass eigenvalues are assumed to be hierarchical, the analysis for the neutrino oscillation experiments requires [1]

$$\frac{2|a|^2}{M_2} \simeq \sqrt{\Delta m^2_{\rm sol}} \simeq (7\times 10^{-5}~{\rm eV}^2)^{1/2}, \quad \frac{2|b|^2}{M_3} \simeq \sqrt{\Delta m^2_{\rm atm}} \simeq (2\times 10^{-3}~{\rm eV}^2)^{1/2}, \tag{8}$$

and $\sin\theta$ should satisfy

$$\sin\theta \simeq \frac{|a|^2/M_2}{\sqrt{2}|b|^2/M_3} \simeq \frac{1}{\sqrt{2}}\sqrt{\frac{\Delta m^2_{\rm sol}}{\Delta m^2_{\rm atm}}} \sim 0.1. \tag{9}$$

This is consistent with the constraint $\sin\theta < 0.16$ which is imposed by the CHOOZ experiment [17]. The effective mass for the neutrinoless double β-decay is estimated in

[3]We adopt these relations in the following study. Under this assumption the model contains seven real parameters. We can loose these strict equalities without changing largely the qualitative results given in this paper.

this model as

$$m_{ee} \lesssim \frac{1}{2}\left|\sqrt{\Delta m^2_{\rm atm}}\sin^2\theta + \sqrt{\Delta m^2_{\rm sol}}\cos^2\theta\right| \sim 2\times 10^{-3}. \tag{10}$$

It seems to be difficult to reach such a value in the next generation experiment.

The lepton flavor violating processes such as $\mu \rightarrow e\gamma$ can constrain the model. It has been suggested that these processes could impose the strong constraint because of the renormalization effect on the soft supersymmetry (SUSY) breaking parameters due to the off-diagonal Yukawa couplings. It can be very severe even in the case of the universal SUSY breaking in the gravity mediation scenario [18]. Here in order to find the conservative condition, we consider the universal soft SUSY breaking in the gravity mediation. The branching ratio of the flavor changing process $\ell_i \rightarrow \ell_j\gamma$ is estimated by taking account of the one-loop contribution as [18]

$$Br(\ell_i \rightarrow \ell_j\gamma) = \frac{\alpha^3\tan^2\beta}{G_F^2 m_{\tilde\ell}^8}\left|\frac{-1}{8\pi^2}(3m_0^2 + A_0^2)(\tilde h^{\nu\dagger}\tilde h^\nu)_{ij}\ln\frac{M_X}{M}\right|^2, \tag{11}$$

where m_0, A_0 and $m_{\tilde\ell}$ represent the soft scalar mass, the SUSY breaking A parameter and the relevant slepton mass, respectively. M_X stands for the unification scale and M is the right-handed neutrino mass scale. Since $\tilde M_1$ is irrelevant to the light neutrino masses as shown in eq. (6), M is appropriate to be taken as M_2 in the present model.

If we assume $m_0 \simeq m_{\tilde\ell} \simeq A_0$ and use eq. (5) for the Yukawa couplings h^ν, each branching ratio for $\mu \rightarrow e\gamma$ and $\tau \rightarrow \mu\gamma$ is estimated as[4]

$$\begin{aligned} Br(\mu \rightarrow e\gamma) &\simeq 3\times 10^{-31}\frac{M_2^2}{m_0^4}\left(\ln\frac{M_X}{M_2}\right)^2\tan^2\beta \leq 1.2\times 10^{-11}, \\ Br(\tau \rightarrow \mu\gamma) &\simeq 4\times 10^{-30}\frac{M_3^2}{m_0^4}\left(\ln\frac{M_X}{M_2}\right)^2\tan^2\beta \leq 1.1\times 10^{-6}. \end{aligned} \tag{12}$$

If we take $M_X = 10^{16}$ GeV and $m_0 = 100$ GeV, for example, the present experimental bounds can be satisfied for $M_2 \lesssim 10^{11}$GeV and $M_3 \lesssim 10^{13}$ GeV even in the case of large $\tan\beta$ such as $\tan\beta \simeq 50$. This means that no lepton flavor violating decays $\ell_i \rightarrow \ell_j\gamma$ contradict with the present model as far as the universal SUSY breaking is assumed even in the gravity mediation scenario.

A remarkable feature of the model is that there are no constraints on M_1 and $\sin\xi$ from the neutrino oscillation data and other present available experiments. If we apply

[4] The decay $\tau \rightarrow e\gamma$ is automatically forbidden as a result of the present texture of the neutrino mass matrix.

this neutrino mass texture to the leptogenesis, M_1 and $\sin\xi$ may be constrained to explain the B asymmetry. In the next section we focus our study on this point.

3 Application to leptogenesis

The decay of the heavy right-handed neutrinos can produce the $B-L$ asymmetry. Then it may explain the B asymmetry in the universe since the sphaleron interaction can convert a part of the $B-L$ asymmetry into the B asymmetry [19]. The L asymmetry or the $B-L$ asymmetry induced through this decay is produced as a result of the CP asymmetry caused by the interference between the tree and one-loop diagrams.

The CP asymmetry appeared in the $\tilde{N}_i$ decay can be generally expressed as [3]

$$\begin{aligned} \varepsilon_i &\equiv \frac{\sum_j \Gamma(\tilde{N}_i \rightarrow L_j H_2) - \sum_j \Gamma(\tilde{N}_i \rightarrow \bar{L}_j \bar{H}_2)}{\sum_j \Gamma(\tilde{N}_i \rightarrow L_j H_2) + \sum_j \Gamma(\tilde{N}_i \rightarrow \bar{L}_j \bar{H}_2)} \\ &= -\frac{1}{8\pi}\frac{1}{(\tilde{h}^\nu \tilde{h}^{\nu\dagger})_{ii}} \sum_{k\neq i} \mathrm{Im}[(\tilde{h}^\nu \tilde{h}^{\nu\dagger})^2_{ik}]\ f\left(\frac{M_k^2}{M_i^2}\right), \end{aligned} \tag{13}$$

where $f(x)$ contains the contributions from both the vertex correction and the self-energy correction, and it has an expression

$$f(x) = \sqrt{x}\left[\ln\left(\frac{1+x}{x}\right) + \frac{2}{x-1}\right]. \tag{14}$$

Applying this formula to our model in which the hierarchical structure for the right-handed neutrino masses is assumed, we obtain

$$\begin{aligned} &\varepsilon_1 \simeq \frac{1}{8\pi}\frac{\sqrt{\Delta m^2_{\rm atm}}\tilde{M}_1}{v^2\sin^2\beta}\sin 2\chi, \qquad \varepsilon_2 \simeq \frac{1}{8\pi}\frac{\sqrt{\Delta m^2_{\rm atm}}M_2}{v^2\sin^2\beta}\sin 2\chi, \\ &\varepsilon_3 \simeq -\frac{1}{16\pi}\frac{\sqrt{\Delta m^2_{\rm sol}}M_2}{v^2\sin^2\beta}\left(\frac{\tilde{M}_1}{M_3}\ln\frac{M_3}{\tilde{M}_1}\sin^2\xi + \frac{M_2}{M_3}\ln\frac{M_3}{M_2}\cos^2\xi\right)\sin 2\chi, \end{aligned} \tag{15}$$

where $\langle H_2\rangle \equiv v\sin\beta$ and $\chi \equiv \arg(a^*b)$. The formulas in eq. (15) show that ε_3 can be much smaller than $\varepsilon_{1,2}$. The interesting point of this result is that ε_1 is almost equal to the expression which saturates its upper bound given in [11].[5] Thus the present texture for the neutrino masses seems to be favorable to induce the B asymmetry.

[5] A model with this feature has been discussed in [14] already. However, it seems not to have a satisfactory structure for the explanation of the neutrino oscillation data.

By using these expression for ε_i, the L asymmetry resulting from the decay of the right-handed neutrinos can be estimated. If we put the excess of the number density of the right-handed neutrino $\tilde{N}_i$ from the equilibrium one as n_i and the entropy density in the comoving volume at the latest $\tilde{N}_i$ decay as s, the produced L asymmetry through this decay can be expressed as

$$Y_L \equiv \frac{n_L}{s} \simeq \sum_{i=1}^{3} \frac{2n_i}{s}\varepsilon_i\kappa_i, \tag{16}$$

where κ_i represents the washout effect which depends on the strength of the Yukawa couplings $\tilde{h}^{\nu}_{ij}$ in eq. (5). The sphaleron interaction which is in the thermal equilibrium at the temperature 10^2 GeV $\lesssim T_{\rm sph} \lesssim 10^{12}$ GeV converts a part of the $B-L$ asymmetry into the B asymmetry in such a way as $n_B/s = -(8/15)(n_L/s)$ in the MSSM case [19, 20]. Thus n_L/s should satisfy $|n_L/s| \gtrsim 10^{-10}$ to realize the observed value $n_B/s \simeq (0.6-1)\times 10^{-10}$.

Since the number density n_i of the right-handed neutrino $\tilde{N}_i$ depends on its generation mechanism, we need to fix it for the quantitative estimation of n_L/s. In the following study we consider both the thermal and non-thermal scenarios for their generation. In the non-thermal leptogenesis we mainly consider that the right-handed neutrinos couple to the inflaton directly and then the inflaton decay produces the right-handed neutrinos non-thermally.

3.1 Thermal leptogenesis

Since the right-handed neutrino masses are supposed to be hierarchical in the present model, the L asymmetry is expected to be produced as a result of the out of equilibrium decay of the lightest right-handed neutrino $\tilde{N}_1$ as studied in a lot of works [5, 6]. In fact, it can be easily checked in the present model. If we use eqs. (5) and (8), the decay width of $\tilde{N}_i$ can be estimated as

$$\begin{aligned}
\Gamma_{\tilde{N}_1} &\simeq \frac{3}{16\pi}\frac{\sqrt{\Delta m^2_{\rm sol}}\tilde{M}_1 M_2}{v^2\sin^2\beta}\sin^2\xi \sim \left(\frac{\tilde{M}_1}{10^7~{\rm GeV}}\right)\left(\frac{M_2}{10^{10}~{\rm GeV}}\right)\sin^2\xi, \\
\Gamma_{\tilde{N}_2} &\simeq \frac{3}{16\pi}\frac{\sqrt{\Delta m^2_{\rm sol}}M_2^2}{v^2\sin^2\beta}\cos^2\xi \sim 10^3\left(\frac{M_2}{10^{10}~{\rm GeV}}\right)^2\cos^2\xi, \\
\Gamma_{\tilde{N}_3} &\simeq \frac{1}{8\pi}\frac{\sqrt{\Delta m^2_{\rm atm}}M_3^2}{v^2\sin^2\beta} \sim 10^9\left(\frac{M_3}{10^{13}~{\rm GeV}}\right)^2,
\end{aligned} \tag{17}$$

where these decay widths are given in the GeV unit. This shows that $\Gamma_{\tilde{N}_1} < \Gamma_{\tilde{N}_2} < \Gamma_{\tilde{N}_3}$ is satisfied. The asymmetry generated by the decay of $\tilde{N}_{2,3}$ is washed out through the L

violating scattering mediated by the thermal $\tilde{N}_1$ *etc.* and then $\kappa_{2,3} \ll 1$.

If we use the thermal number density of the relativistic particle for n_1 and $s = \frac{2\pi^2}{45} g_* T^3$ in eq.(16), the L asymmetry generated through the $\tilde{N}_1$ decay is expressed as

$$\frac{n_L}{s} \simeq \frac{1}{g_*}\varepsilon_1 \kappa_1, \tag{18}$$

where g_* is a degree of freedom for the relativistic particles at this period and $g_* \sim 200$ in the MSSM. In this expression we should note that κ_1 includes also the efficiency factor to generate $\tilde{N}_1$ in the thermal bath other than the washout effect since we suppose that there are no thermal right-handed neutrinos initially.

As is well known in the thermal leptogenesis [6], there is an important quantity $\tilde{m}_1$ which is related to κ_1 and then the strength of the relevant Yukawa couplings of $\tilde{N}_1$. It controls how many $\tilde{N}_1$ is produced in the thermal equilibrium and also how much L asymmetry is washed out. In the present model $\tilde{m}_1$ is expressed as

$$\tilde{m}_1 \equiv (\tilde{h}^\nu \tilde{h}^{\nu\dagger})_{11} \frac{v^2 \sin^2\beta}{\tilde{M}_1} = \left\{ \begin{array}{ll} \frac{3}{2}\sqrt{\Delta m^2_{\rm sol}}\frac{M_2}{M_1}\sin^2\xi & \text{for (a)}, \\ \frac{3}{2}\sqrt{\Delta m^2_{\rm sol}} & \text{for (b)}. \end{array} \right. \tag{19}$$

As is found from eqs. (15) and (19), ε_1 and $\tilde{m}_1$ can be independent from each other because of the freedom of $\sin\xi$. We may expect that there is generally some correlation between these parameters from their definitions (13) and (19). However, the special texture can make them independent in the present model. This feature may cause a substantial influence on the reheating temperature required for the leptogenesis.

If we use the formula in eq. (15), the CP asymmetry ε_1 required from the B asymmetry in the universe is estimated as

$$|\varepsilon_1| \simeq 10^{-8} \times \left(\frac{\tilde{M}_1}{10^8 \text{ GeV}}\right) \gtrsim 10^{-8}, \tag{20}$$

where $g_* \sim 200$ is used. In this estimation the maximum CP phase $|\sin 2\chi| \sim 1$ and $\sin\beta = 1$ are also assumed.[6] From this condition we obtain $\tilde{M}_1 \gtrsim 10^8$ GeV for the lower bound of the $\tilde{N}_1$ mass, which is the ordinary result in the out of equilibrium decay of the thermally produced $\tilde{N}_1$. On the other hand, the effective mass $\tilde{m}_1$ is estimated as

$$\frac{\tilde{m}_1}{10^{-2}\text{ eV}} \simeq \left\{ \begin{array}{ll} \frac{M_2}{M_1}\sin^2\xi & \text{for (a)}, \\ 1 & \text{for (b)}. \end{array} \right. \tag{21}$$

[6]This assumption for $\sin\beta$ brings no crucial difference as far as $\tan\beta$ is in the interesting region such as $1 \lesssim \tan\beta \lesssim 50$.

While $\tilde{m}_1$ takes a fixed value in the case (b), there is the freedom $\sin\xi$ to tune $\tilde{m}_1$ into the desirable value in the case (a). Eq. (21) shows that the efficiency factor can be in the favorable region for the leptogenesis in the case (a) but it seems to be larger than the favorable value in the case (b) [5, 6, 15].

The necessary condition for the out of equilibrium decay of $\tilde{N}_1$ is given by $H > \Gamma_{\tilde{N}_1}$. If we use eq.(17) for this condition, the successful leptogenesis requires that the temperature T at the period of the $\tilde{N}_1$ decay should satisfy

$$T > T_{\rm min} \simeq \sqrt{\tilde{M}_1 M_2}\sin\xi. \tag{22}$$

In the case (b) we find $T_{\rm min} \simeq \tilde{M}_1$ and then $T_R \gtrsim \tilde{M}_1$ should be satisfied. Thus we expect the similar result for the efficiency factor to the previous works [5, 6]. On the other hand, since the case (a) is realized for $M_1 > M_2\sin^2\xi$, we find that $T < \tilde{M}_1$ could be consistent with the condition for the out of equilibrium decay. Such a situation seems to be realized for a sufficiently small $\sin\xi$ without conflicting with the neutrino oscillation data. The small $\sin\xi$ can also make the washout effect negligible. Although this seems to suggest the possibility that the reheating temperature T_R may not be necessary to be high enough compared with $\tilde{M}_1$, however, the sufficient number of $\tilde{N}_1$ may not be produced due to the Boltzmann suppression in the thermal equilibrium. We need to solve a set of Boltzmann equations numerically for the quantitative study of the relation between T_R and $\tilde{M}_1$.

We study these points by solving numerically a set of Boltzmann equations for the MSSM presented in [5]. If we use eq. (8) and assume $|\sin 2\chi| = 1$, the model parameters in this calculation are $M_{1,2,3}$ and $\sin\xi$. As an initial condition for the Boltzmann equations, we assume that both the number density of $\tilde{N}_i$ and the L asymmetry are zero at $z_0 = 0$. A dimensionless parameter z is defined as $z = \tilde{M}_1/T$. The temperature corresponding to z_0 may be considered to correspond to the reheating temperature T_R.

In Fig. 1 we give a solution of the Boltzmann equations with $z_0 = 0.01$ and the input parameters such as $M_1 = 10^9$GeV, $M_2 = 10^{10}$GeV, $M_3 = 10^{13}$GeV and $\sin\xi = 0.02$. This corresponds to the case (a) with $T_R = 10^{11}$GeV. The case (b) cannot yield the sufficient L asymmetry. This can be understood as follows. In this case $\sin\xi$ is required to take a rather large value to realize $M_1 < M_2\sin^2\xi$ satisfying both constraints such as $\tilde{M}_1 \gtrsim 10^8$GeV and $M_3 \lesssim 10^{13}$GeV, which are imposed by the previously discussed phenomenological constraints. Such a $\sin\xi$ makes the $\tilde{N}_1$ Yukawa couplings larger and then the washout effect becomes effective. This is also suggested by eq. (21).

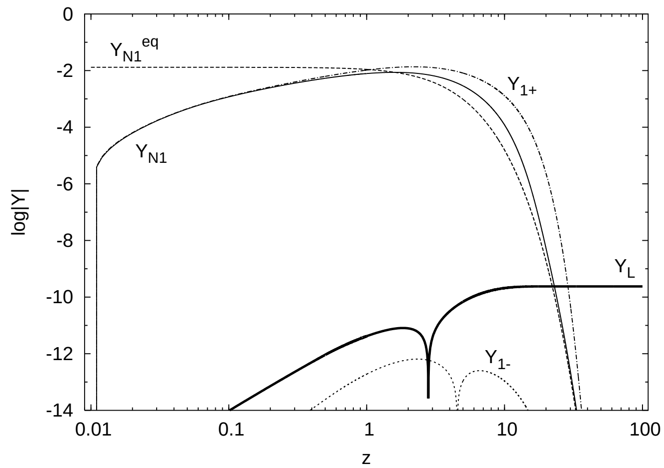

Fig. 1 A typical solution of the Boltzmann equations in the case of the thermal generation of $\tilde{N}_1$. We define Y_i as $Y_i \equiv n_i/s$ $(i = \tilde{N}_1, \tilde{N}_2, L)$ and $Y_{\pm 1} \equiv (n_{SN_1} \pm \bar{n}_{SN_1})/s$ where SN_i stands for the sneutrinos. $Y^{\rm eq}_{N_i}$ is the value in the equilibrium.

In Fig. 2 we show the L asymmetry $|Y_L|$ as a function of $\sin\xi$. In the left panel we plot $|Y_L|$ for the various values of M_1. In the right panel $|Y_L|$ is plotted for the various values of z_0. If we use eqs. (21) and (22), we find that the input parameters adopted to draw Fig. 2 give $\tilde{m}_1 \simeq 10^{-3\sim -4}$eV and $T_R > 10^7$GeV. By using this kind of figures, we can search the lower bounds of M_1 and T_R required to explain the observed B asymmetry for the fixed $M_{2,3}$. We practice this analysis changing the values of $M_{2,3}$ within the allowed region discussed in the previous part. As the result, for the explanation of the B asymmetry based on the present model, we find that the lower bounds of M_1 and T_R can be estimated as

$$M_1 > 5 \times 10^8 \text{ GeV}, \qquad T_R > 6 \times 10^8 \text{ GeV}, \tag{23}$$

and also $\sin\xi$ should be $O(10^{-2})$.

Although the obtained lower bound of the reheating temperature is comparable with the lowest value discussed in other neutrino mass matrix models where $T_R \simeq 10^{9-10}$GeV is usually suggested, we cannot make it much lower. Since the lower bound of the mass eigenvalue $\tilde{M}_1$ is determined by eq. (20), this result seems not to be avoided in the thermal leptogenesis as far as we do not assume the degeneracy among the masses of the right-handed neutrinos. Recently, in [21] the upper bound for the reheating temperature required from the gravitino problem is estimated as 10^{5-7}GeV if the gravitino has the mass in the range 10^{2-3}GeV. If we do not suppose the light gravitino scenario and we

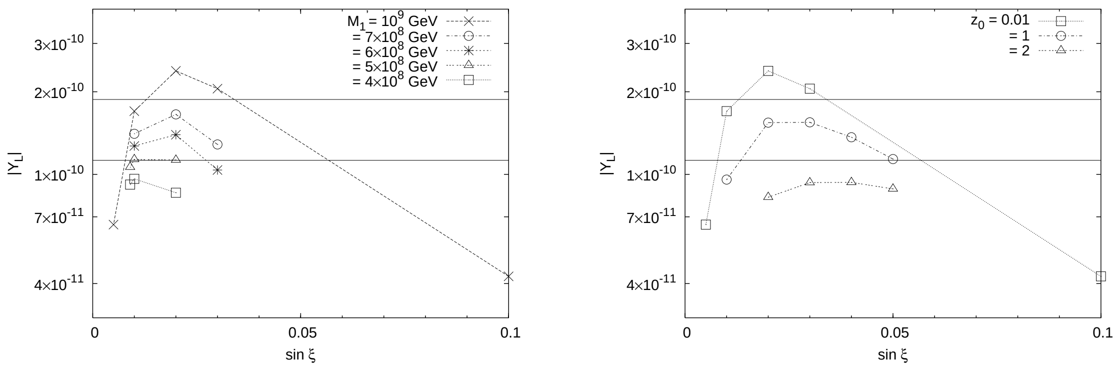

Fig. 2 The L asymmetry $|Y_L|$ as a function of $\sin\xi$. Horizontal thin lines represent the desirable region to explain the observed B asymmetry. In the left panel M_1 is varied keeping other parameters fixed in such a way as $M_2 = 10^{10}$GeV, $M_3 = 10^{13}$GeV and $z_0 = 0.01$. In the right panel we vary the z_0 value keeping others fixed as $M_1 = 10^9$GeV, $M_2 = 10^{10}$GeV and $M_3 = 10^{13}$GeV.

follow this bound, the present model is unable to be reconciled with the gravitino problem. We need to consider the initial $\tilde{N}_1$ to be yielded in other way. As such a possibility, we study the non-thermal leptogenesis in the next part.

3.2 Nonthermal leptogenesis

In this subsection we consider that the right-handed neutrinos $\tilde{N}_i$ are produced through the decay of the inflaton. This kind of model has been discussed in [9]. The interaction between the inflaton superfield Φ and $\tilde{N}_i$ is assumed to be given by the superpotential

$$W = \sum_{i=1}^{3} \lambda_i \Phi \tilde{N}_i^2. \tag{24}$$

After the inflation ends, the inflaton ϕ starts to oscillate and decays to reheat the universe into the temperature T_R. A part of its oscillation energy ρ of the inflaton is converted into $\tilde{N}_i$ through its decay at $H \simeq \Gamma_\phi$. The decay width Γ_ϕ of the inflaton ϕ can be expressed as

$$\Gamma_\phi = \sum_{i=1}^{3} \frac{\lambda_i^2}{4\pi} m_\phi + \cdots, \tag{25}$$

where the ellipses stand for the contribution from other decay modes of the inflaton and we assume them to be negligible. The coupling constants λ_i are constrained in such a way as

$$\sum_{i=1}^{3} \lambda_i^2 \simeq 10^{-21} \left(\frac{10^{16}\ \mathrm{GeV}}{m_\phi}\right) \left(\frac{T_R}{10^6\ \mathrm{GeV}}\right)^2, \tag{26}$$

which is derived from the condition $H \simeq \Gamma_\phi$. From this we find that the couplings between the right-handed neutrinos and the inflaton can be small enough not to affect the inflaton potential. The inflaton mass m_ϕ can depend on the assumed inflation model. However, it should satisfy $m_\phi > \tilde{M}_i$ to guarantee the inflaton decay into the right-handed neutrinos $\tilde{N}_i$.

If we use B_i to denote the branching ratio for the decay $\phi \to N_i^2$, we have the energy relation $\rho B_i = \tilde{M}_i n_i$ where $\rho = \frac{\pi^2}{30} g_* T_R^4$. Thus the non-thermally generated number density n_i of $\tilde{N}_i$ can be written as

$$\frac{n_i}{s} = \frac{3 T_R B_i}{4 \tilde{M}_i}. \tag{27}$$

As we mentioned below eq. (15), ε_3 is much smaller than $\varepsilon_{1,2}$. In the case dominated by B_3 the produced L asymmetry is expected to be very small unless T_R is rather large. Then we concentrate our study into the case dominated by $B_{1,2}$. Since we find $\varepsilon_1/\tilde{M}_1 = \varepsilon_2/M_2$ from eq. (15), the L asymmetry generated through the immediate decay of $\tilde{N}_{1,2}$ is estimated by using eq. (16) as

$$\begin{aligned} \frac{n_L}{s} &\simeq \frac{3}{16\pi} \frac{\sqrt{\Delta m^2_{\rm atm}} T_R}{v^2 \sin^2\beta} (\kappa_1 B_1 + \kappa_2 B_2) \sin^2 2\chi \\ &\simeq 10^{-10} \left(\frac{T_R}{10^6 \text{ GeV}} \right) (\kappa_1 B_1 + \kappa_2 B_2) \sin^2 2\chi. \end{aligned} \tag{28}$$

Eq. (28) suggests that T_R should be larger than 10^6 GeV to explain the B asymmetry in any case. Moreover, we can find that the $\tilde{M}_i$ dependence of n_L/s is confined into the washout factor κ_i. If we consider that $\tilde{M}_i$ is larger than T_R, the washout effect due to $\tilde{N}_i$ is expected to be suppressed by the Boltzmann factor. Thus for the larger $\tilde{M}_i$ in such a region, n_L/s becomes insensitive for the change of the values of $\tilde{M}_i$.

The estimation of the L asymmetry in eq. (28) is justified only if $\tilde{N}_i$ decays into the light fields immediately after its production [9]. This requires that $H \simeq \Gamma_\phi \lesssim \Gamma_{\tilde{N}_i}$ should be satisfied for all $\tilde{N}_i$ which have the substantial branching ratio B_i.[7] Since the inflaton decay width can be estimated by using $H \simeq \Gamma_\phi$ as

$$\Gamma_\phi \simeq 0.3 g_*^{1/2} \frac{T_R^2}{M_{\rm pl}}, \tag{29}$$

[7] In the construction of the mass matrices presented in appendix A, we find that $B_2 \gg B_1$ is satisfied if the inflaton has no global charges. However, we consider the general case here.

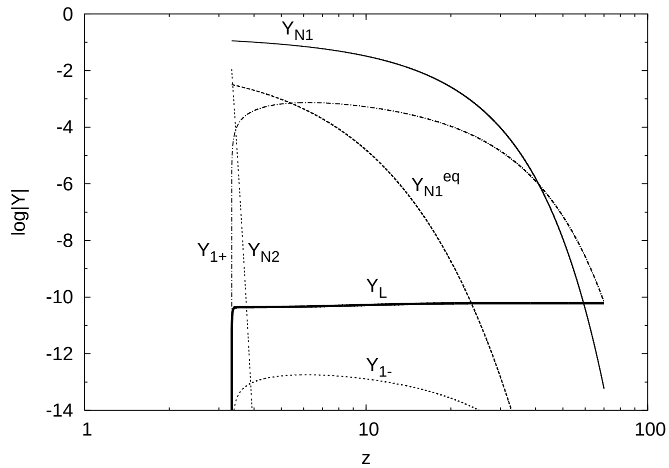

Fig. 3 A typical solution of the Boltzmann equations in the case of the non-thermal generation of $\tilde{N}_i$. The definitions of Y_i are the same as the ones in Fig. 1.

we can write the condition for the justification of eq. (28) by applying eqs. (26) and (29) to $\Gamma_\phi \lesssim \Gamma_{\tilde{N}_1}$ in the form as

$$\left(\frac{T_R}{10^6\ \mathrm{GeV}}\right)^2 \lesssim 10^6 \left(\frac{\tilde{M}_1}{10^7\ \mathrm{GeV}}\right)\left(\frac{M_2}{10^{10}\ \mathrm{GeV}}\right)\sin^2\xi. \tag{30}$$

Since this condition can be easily satisfied for the desirable values of T_R, $\tilde{M}_1$ and M_2, we find that eq. (28) can be validated in our interested case. However, it is also possible that the immediate decay condition is not satisfied for $\tilde{N}_1$. This occurs for $\Gamma_{\tilde{N}_1} < \Gamma_\phi < \Gamma_{\tilde{N}_2}$ in the case of $B_1 \simeq B_2$. In that case we should take account that the L asymmetry produced through the $\tilde{N}_2$ decay may be washed out by the late entropy release due to the $\tilde{N}_1$ decay other than by the usual thermal washout. This effect is discussed in appendix B. We should also check that the condition (30) can be consistent with the above mentioned condition $m_\phi > M_{\tilde{N}_i}$. This consistency can be easily checked by applying eq. (26) to Γ_ϕ.

In order to study the relation between T_R and $\tilde{M}_i$, it is useful to classify the situation into three cases: (i) $T_R \lesssim \tilde{M}_1$, (ii) $\tilde{M}_1 \lesssim T_R \lesssim M_2$, (iii) $M_2 \lesssim T_R$. Among these three cases, both the cases (i) and (ii) can easily satisfy the condition (30). On the other hand, the case (iii) satisfies it only for $M_2 \lesssim 10\tilde{M}_1 \sin^2\xi$, which requires $\sin\xi \gg 0.1$. Thus only the two cases (i) and (ii) seem to be promising for the leptogenesis consistent with the low reheating temperature. In fact, $T_R \simeq 10^6$ GeV may be allowed in these cases with $B_{1,2} = O(1)$ and $|\sin 2\chi| = O(1)$ if κ_1 or κ_2 can be $O(1)$.

The washout effect is expected to be mainly caused by the L violating interactions

due to the thermal $\tilde{N}_i$. Since there is the Boltzmann suppression for these processes in the case (i), their substantial washout effect cannot be expected. On the other hand, in the case (ii) $\tilde{N}_1$ can contribute to the washout of the L asymmetry since there is no large Boltzmann suppression. To escape this situation the smallness of the $\tilde{N}_1$ Yukawa couplings is required. This may be realized for the small $\sin\xi$ case.

To take account of the washout effect quantitatively, we need to solve the Boltzmann equations numerically by using eq. (27) as the initial value for the n_i at $z_0 = M_1/T_R$. In Fig. 3 we show a typical solution for the Boltzmann equations in the case (a). In this figure we assume $|\sin 2\chi| = 1$, $B_1 = B_2 = 0.5$ and $T_R = 3\times 10^6$GeV. The input parameters are taken as $M_1 = 10^7$GeV, $M_2 = 10^8$GeV, $M_3 = 10^{13}$GeV and $\sin\xi = 0.01$. This figure shows that the number density of $\tilde{N}_2$ rapidly decreases following the Boltzmann distribution. The L asymmetry reaches the final value faster compared with the thermal case. In the case (b) the sufficient amount of the L asymmetry cannot be produced. The reason is considered to be the same as the thermal case.

In Fig. 4 we plot the L asymmetry $|Y_L|$ as the function of $\sin\xi$ for various values of the input parameters. We calculate Y_L for the typical three models with different branching ratios and plot them with the different symbols, that is, the squares for $B_1 = 0$, $B_2 = 1$, the circles for $B_1 = B_2 = 0.5$ and the triangles for $B_1 = 1$, $B_2 = 0$. The reheating temperature is assumed to be $T_R = 3\times 10^6$GeV. The left panel corresponds to the case (i). In this figure, as the input parameters we use $M_1 = 10^8$GeV, $M_2 = 10^9$GeV, $M_3 = 10^{13}$GeV for three types of the black symbols and $M_1 = 10^8$GeV, $M_2 = 5\times 10^8$GeV, $M_3 = 10^{13}$GeV for the white symbols. The right panel corresponds to the case (ii). In this case, as the input parameters we use $M_1 = 10^5$GeV, $M_2 = 5\times 10^8$GeV, $M_3 = 10^{13}$GeV for the black symbols and $M_1 = 10^5$GeV, $M_2 = 10^8$GeV, $M_3 = 10^{13}$GeV for the white symbols. The typical feature in these cases is that the $\sin\xi$ value can be smaller compared with the thermal case since we need not to produce $\tilde{N}_i$ thermally.

In both panels of Fig. 4 the larger $|Y_L|$ is realized for the larger B_2 since the washout effect is smaller compared with the smaller B_2 case. If we make M_2 larger keeping M_1 fixed in these figures, $|Y_L|$ becomes a little bit larger but it seems to reach almost the upper bound in this setting. In the left panel the condition (30) is satisfied only for $\sin\xi \gtrsim 10^{-3}$. In the case of $B_2 = 1$, however, this condition should be replaced by $\Gamma_\phi \lesssim \Gamma_{\tilde{N}_2}$ and it is satisfied for all region of $\sin\xi$ in this figure. In the case of $B_1 = B_2 = 0.5$, for $\sin\xi < 10^{-3}$

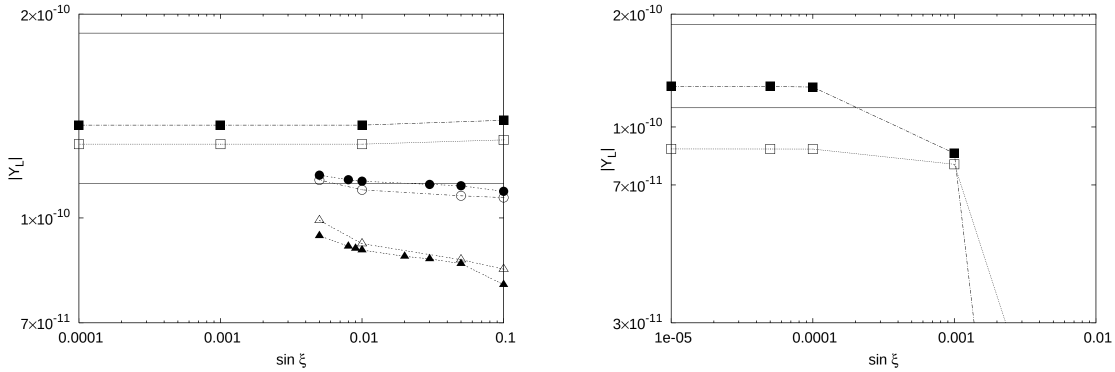

Fig. 4 The L asymmetry Y_L as a function of $\sin\xi$. Horizontal thin lines represent the desirable region to explain the observed B asymmetry. The explanations for the symbols are presented in the text.

we should take account of the additional washout effect to the L asymmetry produced through the $\tilde{N}_2$ decay, which is discussed in appendix B. It introduces the suppression factor $\sqrt{\tilde{M}_1 M_2}\sin\xi/T_R$. It takes a value smaller than $O(10^{-1})$ in the present case. On the other hand, the L asymmetry produced by the late $\tilde{N}_1$ decay cannot be a sufficient amount because of the low reheating temperature. Thus we find that the models with the substantial B_1 cannot explain the B asymmetry for $\sin\xi < 10^{-3}$. The Yukawa couplings of $\tilde{N}_1$ and $\tilde{N}_2$ are proportional to $\sin\xi$ and $\cos\xi$, respectively. This fact affects the behavior of $|Y_L|$ as the function $\sin\xi$. In fact, the figures show the slight increase in the case of $B_2 = 1$ and the decrease in the case of $B_1 = 1$ when the $\sin\xi$ value becomes larger.

In the right panel the condition (30) is satisfied only for $\sin\xi > 0.1$. The situation is the same as the left panel in the case $B_2 = 1$. Thus in the case (ii) only the model with $B_2 \sim 1$ can have the possibility to explain the B asymmetry. The magnitude of $|Y_L|$ becomes smaller for the larger $\sin\xi$ even in the case of $B_2 = 1$. Since $\tilde{N}_1$ can be produced thermally, the large $\sin\xi$ makes the washout effect more effective, This feature explains the $|Y_L|$ behavior against $\sin\xi$ in the right panel.

We practice this kind of analysis changing the input parameters within the allowed region. As a result of such study, we find that the lower bound of the required reheating temperature to explain the B asymmetry is $T_R \simeq 3 \times 10^6$ GeV for both cases (i) and (ii).

Finally, we briefly comment on other possibility for the non-thermal leptogenesis [10]. In the early universe the scalar potential of the sneutrino $\tilde{N}_1$ may be flat enough to deviate largely from its potential minimum.[8] If this happens during the inflation, the condensate

[8]The sneutrino $\tilde{N}_1$ can be an inflaton itself as discussed in [22]. However, we do not consider this

of $\tilde{N}_1$ starts to oscillate at $H \simeq \tilde{M}_1$ and decays at $H \simeq \Gamma_{\tilde{N}_1}$.[9] This oscillation may dominate the energy density of the universe at a certain time after the reheating due to the inflaton decay because of its behavior as a matter.[10] We assume that it is the case here.

Since its energy density is expressed by $\rho_{\tilde{N}_1} = \tilde{M}_1^2|\tilde{N}_1|^2$, the $\tilde{N}_1$ number density n_1 is estimated as $\tilde{M}_1|\tilde{N}_1|^2$. Thus the ratio of the L asymmetry produced through its decay to the entropy density is estimated as

$$\frac{n_L}{s} = \frac{2\tilde{M}_1|\tilde{N}_1|^2}{s}\varepsilon_1\kappa_1 = \frac{3}{2}\frac{T_R}{\tilde{M}_1}\varepsilon_1\kappa_1, \tag{31}$$

where in the last equality we use the above mentioned assumption $\rho_{\tilde{N}_1} = \frac{\pi^2}{30}g_* T_R^4$ for the energy density. If we use eq. (15) for the CP asymmetry ε_1, we obtain the similar result for n_L/s to the previous example. However, in this case $\Gamma_\phi > \Gamma_{\tilde{N}_1}$ should be satisfied and this condition imposes $T_R > \sqrt{10\tilde{M}_1 M_2}\sin\xi$. Thus the expected L asymmetry is estimated as

$$\frac{n_L}{s} \simeq 10^{-10}\frac{T_R}{10^6\ \mathrm{GeV}} \gtrsim 10^{-10}\frac{\sqrt{\tilde{M}_1 M_2}\sin\xi}{10^5\ \mathrm{GeV}}, \tag{32}$$

where we assume $|\sin 2\chi| = 1$ and $\kappa_1 = 1$. This relation gives a constraint for the undetermined parameters in the present neutrino mass texture based on the B asymmetry. The condition $\Gamma_\phi > \Gamma_{\tilde{N}_1}$ also requires us to consider the different type of the inflation from the previous non-thermal example.

Since $\kappa_1 \simeq 1$ is validated only for the case $T_R < \tilde{M}_1$, eq. (32) cannot be applied to the case (b) in which $T_R \gtrsim \tilde{M}_1$ follows. We can obtain a sufficient amount of the L asymmetry by taking $\tilde{M}_1$ large enough to make ε_1 large but keeping $\sin\xi$ small enough to realize $T_R < \tilde{M}_1$. Thus, in the case (a) a low reheating temperature like $T_R \simeq 10^6$GeV can be enough to produce the required B asymmetry by setting M_1 and $\sin\xi$ suitably. To realize such a low reheating temperature, for example, the mass parameters in the neutrino sector may be taken as

$$M_1 = 10^8\ \mathrm{GeV}, \quad M_2 = 10^{10}\ \mathrm{GeV}, \quad M_3 = 10^{13}\ \mathrm{GeV}, \quad \sin\xi = 10^{-4}. \tag{33}$$

possibility since the reheating temperature is too high to be reconciled with the gravitino problem in this case.

[9] This oscillation may start during the inflation ($\Gamma_\phi < \tilde{M}_1$) or after the inflation ($\Gamma_\phi > \tilde{M}_1$).

[10] During this oscillation, the flat direction may store the L asymmetry due to the Affleck-Dine mechanism as discussed in [23]. However, we do not consider this possibility here.

Since the effective mass $\tilde{m}_1$ is estimated as $\tilde{m}_1 \sim 10^{-8}$eV, the washout effect is completely negligible as expected. We have no gravitino problem in this case since the reheating temperature realized by the decay of the $\tilde{N}_1$ condensate is sufficiently low comparable to the one given in [21].

4 Summary

We have proposed the neutrino mass matrices in the framework of the MSSM extended with the three generation right-handed neutrinos. These mass matrices can realize the bi-large mixing among the neutrino flavors and explain the neutrino oscillation data. It can also saturate the upper bound of the CP asymmetry ε_1 appeared in the leptogenesis. Although this model is composed of rather restricted number of parameters, it can make the CP asymmetry ε_1 and the effective neutrino mass $\tilde{m}_1$ independent. We have applied this model to the thermal and non-thermal leptogenesis and studied the influence of this feature on the reheating temperature, which is crucial for the cosmological gravitino problem.

In the thermal leptogenesis our neutrino mass texture seems not to be able to make the reheating temperature required from the explanation of the B asymmetry low enough to be consistent with the gravitino problem. However, it seems to be able to realize a value near the lower bound of the reheating temperature obtained in the thermal leptogenesis framework. In the non-thermal case we have found that the low reheating temperature consistent with the gravitino problem can be sufficient for the successful leptogenesis. Even in that case the parameters in the neutrino mass matrices can be consistent with the neutrino oscillation data.

As is shown in this study, some kinds of neutrino mass texture can be constrained by the leptogenesis. It may be worthy to proceed a lot of study based on the concrete neutrino model to clarify the relation among the neutrino mass texture, the leptogenesis and the gravitino problem.

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Appendix A

In this appendix we present an example of the construction for the assumed texture of the neutrino mass matrix. We consider the Frogatt-Nielsen type global flavor symmetry $\mathrm{U}(1)^5$ and an additional discrete Z_2 symmetry in the lepton sector. The charge assignment of the chiral superfields for the symmetry $\mathrm{U}(1)^5 \times Z_2$ is assumed as

$$\begin{aligned}
&N_1\ (1,1,1,0,1;+), \quad N_2\ (1,0,1,1,1;+), \quad N_3\ (1,1,0,1,0;+), \\
&L_1\ (0,0,-1,-1,0;+), \quad L_2\ (-1,0,0,-1,0;+), \quad L_3\ (0,-1,0,-1,0;+).
\end{aligned} \tag{34}$$

The Higgs chiral superfield H_2 is neutral for this symmetry. In order to realize the hierarchical structure of the mass matrices, we introduce the following several chiral superfields which are singlet for the standard gauge groups:

$$\begin{array}{ccc}
\phi_1\ (-1,-1,-1,0,0;-) & \phi_2\ (-1,0,-1,-1,0;-), & \phi_3\ (-2,-2,0,-2,0;+), \\
\chi_1\ (-1,0,0,0,0;+), & \chi_2\ (0,-1,0,0,0;+), & \chi_3\ (0,0,-1,0,0;+), \\
 & \eta\ (0,0,0,0,-1;+). &
\end{array} \tag{35}$$

If we assume that the scalar components of these superfields get vacuum expectation values defined by

$$\epsilon_i \equiv \frac{\langle\phi_i\rangle}{M_{\rm pl}}, \qquad \delta \equiv \frac{\langle\chi_i\rangle}{M_{\rm pl}}, \qquad \zeta \equiv \frac{\langle\eta\rangle}{M_{\rm pl}}, \tag{36}$$

we can obtain both the right-handed Majorana mass matrix and the Dirac mass matrix as follows:

$$\mathcal{M} \simeq M_3 \begin{pmatrix} \epsilon_1^2\zeta^2/\epsilon_3 & \epsilon_1\epsilon_2\zeta^2/\epsilon_3 & 0 \\ \epsilon_1\epsilon_2\zeta^2/\epsilon_3 & \epsilon_2^2\zeta^2/\epsilon_3 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \qquad m_D = v_2 \begin{pmatrix} 0 & 0 & 0 \\ \zeta\delta & \zeta\delta & 0 \\ 0 & \delta & \delta \end{pmatrix}, \tag{37}$$

where $M_3 \equiv M_{\rm pl}\epsilon_3$ and the order one coefficients are abbreviated. The difference between the cases (a) and (b) should be considered to be explained by these coefficients.

We can check that these mass matrices can consistently realize the texture assumed in the text. Comparing m_D in eq. (2) with that in eq. (37) and also using eq. (8), we find

$$\delta \sim 0.1 \left(\frac{M_3}{10^{13}\mathrm{GeV}}\right)^{1/2}, \qquad \zeta \sim 0.4 \left(\frac{M_2}{M_3}\right)^{1/2}, \qquad \epsilon_3 = \frac{M_3}{M_{\rm pl}}. \tag{38}$$

Applying this result to $\mathcal{M}$ in eq. (2) and eq. (37), we obtain

$$M_1 \sim \left(\frac{\epsilon_1}{\epsilon_2}\right)^2 M_2, \qquad \epsilon_3 \sim 0.16\epsilon_2^2, \qquad \sin\xi \sim \frac{\epsilon_1}{\epsilon_2}. \tag{39}$$

If we take $\delta \sim 0.1$, $\zeta \sim 0.03$ and $\epsilon_2 \simeq 30\epsilon_1$, for example, we find

$$M_1 \sim 10^8 \text{ GeV}, \quad M_2 \sim 10^{10} \text{ GeV}, \quad M_3 \sim 10^{13} \text{ GeV}, \quad \sin\xi \sim 0.03. \tag{40}$$

This suggests that (37) can realize the mass matrices assumed in the text.

We can consider the couplings of the right-handed neutrinos $\tilde{N}_i$ to the inflaton ϕ, which is assumed to be neutral under the present global symmetry. The lowest order superpotential allowed by this symmetry is written as

$$W = c_1 \frac{\phi_1^2 \eta^2}{M_{\rm pl}^4} \phi N_1^2 + c_2 \frac{\phi_2^2 \eta^2}{M_{\rm pl}^4} \phi N_2^2 + c_3 \frac{\phi_3}{M_{\rm pl}} \phi N_3^2, \tag{41}$$

where the coefficients $c_{1,2,3}$ are assumed to be O(1). Thus the coupling constants λ_i are estimated as

$$\lambda_1 = c_1 \epsilon_1^2 \zeta^2, \qquad \lambda_2 = c_2 \epsilon_2^2 \zeta^2, \qquad \lambda_3 = c_3 \epsilon_3. \tag{42}$$

We find that these couplings can be consistent with the low value of T_R if $m_\phi \simeq 10^{12}$GeV is assumed. In this case $B_2 \gg B_1$ can be satisfied since $B_3 = 0$ is realized for $M_3 \simeq 10^{13}$GeV.

Appendix B

In the non-thermal leptogenesis the L asymmetry produced through the $\tilde{N}_2$ decay may be washed out in a different way compared with the thermal case. If $\Gamma_{\tilde{N}_1} < \Gamma_\phi \lesssim \Gamma_{\tilde{N}_2}$ is satisfied, $\tilde{N}_2$ decays into the light particles at the time t_2 immediately after the inflaton decays into $\tilde{N}_2$. The decay products of $\tilde{N}_2$ behaves as the radiation and its energy density decreases as $\rho_{\tilde{N}_2} \propto a^{-4}$ where a is the cosmological scale parameter. On the other hand, $\tilde{N}_1$ decays at the time t_1 after the completion of the inflaton decay. If $\tilde{N}_1$ behaves as a matter because of $T_R < \tilde{M}_1$, its energy density decreases as $\rho_\phi \propto a^{-3}$. Then the cosmological energy density may be dominated by the $\tilde{N}_1$ energy at least as far as B_1 and B_2 are the same order. The additional washout can occur in such a case.

The cosmological energy density $\rho(t_2)$ and the temperature $T_2(t_2)$ of its decay products can be expressed as

$$\rho(t_2) B_2 = \frac{\pi^2}{30} g_*(t_2) T_2^4(t_2), \quad H^2(t_2) = \frac{\rho(t_2)}{3M_{\rm pl}^2} \simeq \Gamma_\phi^2. \tag{43}$$

Taking account of these, we can find that the temperature T_2 of the decay products of $\tilde{N}_2$ satisfies the relation such as

$$\left(\frac{T_2(t_2)}{T_2(t_1)}\right)^4 = \left(\frac{a(t_1)}{a(t_2)}\right)^4 = \left(\frac{H(t_2)}{H(t_1)}\right)^{8/3} = \left(\frac{\Gamma_\phi}{\Gamma_{\tilde{N}_1}}\right)^{8/3}. \tag{44}$$

From this relation we obtain

$$T_2(t_2) = T_2(t_1)\left(\frac{\Gamma_\phi}{\Gamma_{\tilde{N}_1}}\right)^{2/3}. \tag{45}$$

Now we can estimate the entropy production of the late decay of $\tilde{N}_1$. Since $H^2(t_1) = \rho(t_1)/3M_{\rm pl}^2 = \Gamma_{\tilde{N}_1}$ is satisfied, the ratio between the entropy density $s_b(t_1)$ before the $\tilde{N}_1$ decay and the entropy density $s_a(t_1)$ after the decay can be written as

$$\frac{s_b(t_1)}{s_a(t_1)} = \frac{g_*(t_2)T_2^3(t_1)}{g_*(t_1)T_1^3(t_1)} \simeq \left(\frac{T_2(t_2)}{T_1(t_1)}\right)^3\left(\frac{\Gamma_{\tilde{N}_1}}{\Gamma_\phi}\right)^2 \simeq \left(\frac{\Gamma_{\tilde{N}_1}}{\Gamma_\phi}\right)^{1/2}. \tag{46}$$

By using the expression of each decay width, we find that κ_2 is written as

$$\kappa_2 \simeq \kappa\frac{\sqrt{\tilde{M}_1 M_2}}{T_R}\sin\xi, \tag{47}$$

where κ is the usual thermal washout effect due to the L violating scattering mediated by the right-handed neutrinos and so on.

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